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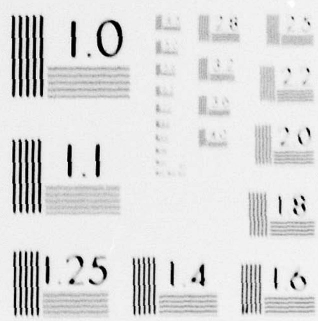


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RADC-TR-79-96
In-House Report
March 1979

# SPECKLE INTERFEROMETRY FOR A PARTIALLY COHERENT SOURCE

Ronald L. Fante

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Chief, Antennas and RF Components Branch

APPROVED: Allan C. Schell

ALLAN C. SCHELL
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## REPORT DOCUMENTATION PAGE

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| | |
|---|---|
| 1. REPORT NUMBER: (14) RADC-TR-79-96 | 2. GOVT ACCESSION NO. |
| | 3. RECIPIENT'S CATALOG NUMBER |
| 4. TITLE (and Subtitle): (6) SPECKLE INTERFEROMETRY FOR A PARTIALLY COHERENT SOURCE | 5. TYPE OF REPORT & PERIOD COVERED: In-House Report |
| | 6. PERFORMING ORG. REPORT NUMBER |
| 7. AUTHOR(s): (10) Ronald L. Fante | 8. CONTRACT OR GRANT NUMBER(s) |
| 9. PERFORMING ORGANIZATION NAME AND ADDRESS: Deputy for Electronic Technology (RADC/EEA) Hanscom AFB Massachusetts 01731 | 10. PROGRAM ELEMENT, PROJECT, TASK AREA & WORK UNIT NUMBERS: 61102F (16) 2305J303 (17) J3 |
| 11. CONTROLLING OFFICE NAME AND ADDRESS: Deputy for Electronic Technology (RADC/EEA) Hanscom AFB Massachusetts 01731 (11) | 12. REPORT DATE: Mar[redacted]979 |
| | 13. NUMBER OF PAGES: 15 |
| 14. MONITORING AGENCY NAME & ADDRESS (if different from Controlling Office): (12) 16p. | 15. SECURITY CLASS. (of this report): Unclassified |
| | 15a. DECLASSIFICATION/DOWNGRADING SCHEDULE |

16. DISTRIBUTION STATEMENT (of this Report)

Approved for public release; distribution unlimited.

17. DISTRIBUTION STATEMENT (of the abstract entered in Block 20, if different from Report)



18. SUPPLEMENTARY NOTES

19. KEY WORDS (Continue on reverse side if necessary and identify by block number)

Turbulence
Partial coherence
Speckle

20. ABSTRACT (Continue on reverse side if necessary and identify by block number)

We have obtained the spatial frequency transfer function for speckle interferometry for the case when the source is partially coherent (spatially) and can be approximated by the quasi-stationary model.

The transfer function has been examined, and relatively simple approximations to it have been found for the limiting cases of low and high spatial frequencies. These results reduce to previous results in the limiting case when the source is completely incoherent.

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## Contents



## Illustrations

# Speckle Interferometry
# For a Partially Coherent Source

## 1. INTRODUCTION

The conventional theoretical result[1-9] for the spatial frequency spectrum of a source, as obtained using speckle interferometry[4, 7] is based on the assumption that the source is spatially incoherent. This assumption is not restrictive when imaging distant stars, because the coherence patches are generally unresolvable by the imaging system; however, when one tries to obtain detail over sizes on the order of the spatial coherence length, the conventional theory is no longer valid, and must be modified. In this paper we will obtain the frequency transfer function for the combination of atmosphere and imaging system in the case when the source is partially coherent, and can be represented by the quasi-stationary model employed by Carter and Wolf,[10] and Leader.[11] The limiting case when the source consists of 2 point sources with random phase relationship has been studied by Miller and Korff.[12]

(Received for publication 29 March 1979)

Due to the number of references to be included as footnotes on this page, the reader is referred to the list of references, page 15.



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## 2. ANALYTICAL FORMALISM

In the Labeyrie[6, 7] procedure a series of short exposure photographs of a star are taken and then successively projected, in the conventional manner, by a laser onto a second piece of film. This leads to an evaluation of the ensemble average of the modulus squared of the Fourier transform of the irradiance, and gives finer detail than conventional seeing calculations predict should be resolvable in the presence of atmospheric turbulence.

In order to discuss the aforementioned method quantitatively, let us refer to Figure 1. If the field distribution of the source is denoted by $u_o(\underline{\rho}_1)$, the field in the image plane, where film is located is,

$$u_f(\underline{\rho}) = \frac{e^{-ik(L+x)}}{\lambda^2 Lx} \iint_{-\infty}^{\infty} d^2\rho_1\, u_o(\underline{\rho}_1)\, M(\underline{\rho}_1, \underline{\rho}) \tag{1}$$

where

$$M(\underline{\rho}_1, \underline{\rho}) = \iint_{-\infty}^{\infty} d^2\rho_2\, T(\underline{\rho}_2) \exp\left\{ ik\underline{\rho}_2 \cdot \left(\frac{\underline{\rho}_1}{x} + \frac{\underline{\rho}}{L}\right) + \psi(\underline{\rho}_2, \underline{\rho}_1) \right\} \quad . \tag{2}$$

In writing Eq. (1) we have assumed that* the film is in the Fraunhofer zone of the source, defined $T(\underline{\rho}_2)$ as the transmissivity of the aperture and $k = 2\pi/\lambda$ as the signal wavenumber. Also $\psi(\underline{\rho}_2, \underline{\rho}_1)$ is the additional complex phase, due to atmospheric turbulence, of a spherical wave propagating from the point $(x, \underline{\rho}_1)$ on the source to the point $(0, \underline{\rho}_2)$ in the aperture. The short-exposure** transparency on the film is proportional to $u_f u_f^*$. N such short exposures are made, with the time interval between successive exposures greater than the atmospheric coherence time, so that we have N statistically independent short exposure photographs. If a photographic plate is then placed at $x = -L$, and these exposures successively projected onto it, it is found that the transparency on the plate is proportional to $\langle |\hat{I}_f(k\underline{\rho}/L)|^2 \rangle$, where

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*If the film is in the Fresnel zone, we simply replace $u_f(\underline{\rho})$ by $u_f(\underline{\rho}) \exp(ik\rho^2/2L)$, and $u_o(\underline{\rho}_1)$ by $u_o(\underline{\rho}_1) \exp(-ik\rho_1^2/2x)$.

**By short exposure we mean short compared with the coherence time of the atmosphere.

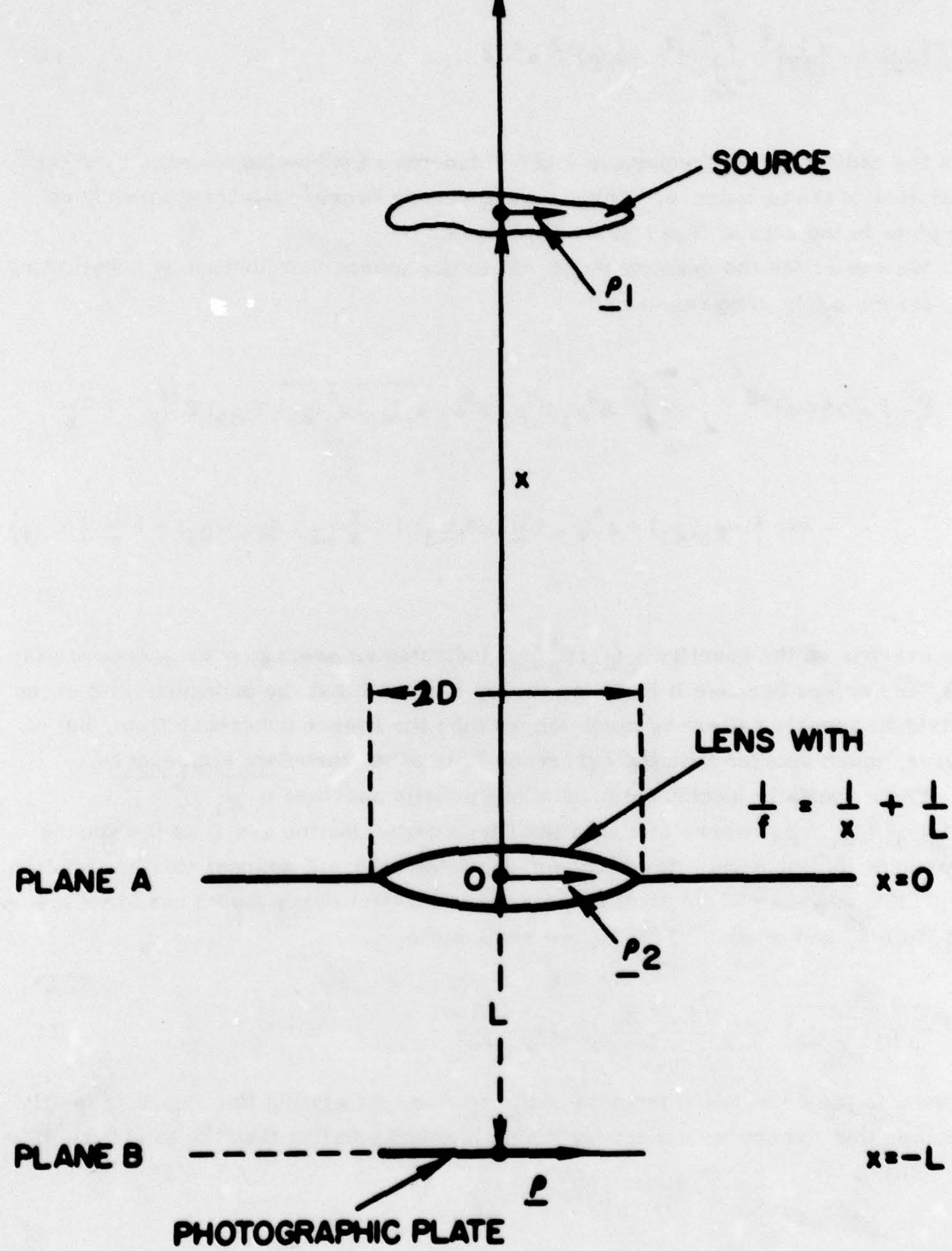


Figure 1. Geometry Assumed for the Imaging System

$$\hat{I}_f(\underline{\omega}) = \left(\frac{1}{2\pi}\right)^2 \iint_{-\infty}^{\infty} d^2\rho \, \langle |u_f(\underline{\rho})|^2 \rangle \, e^{i\underline{\omega}\cdot\underline{\rho}} \,, \tag{3}$$

$\underline{\omega}$ is the radian spatial frequency, and $\langle\ \rangle$ denotes an ensemble average over the statistics of the turbulence. This average occurs because the transparency on the plate is the sum of $N \gg 1$ short exposures.

We can relate the quantity in Eq. (3) to the source distribution by substituting Eq. (1) for $u_f(\underline{\rho})$. The result is

$$\hat{I}_f(\underline{\omega}) = (2\pi\lambda x)^{-2} \int \cdots \int_{-\infty}^{\infty} d^2\rho_1 \, d^2\rho_2 \, d^2\rho_3 \, \overline{u_o(\underline{\rho}_1) u_o^*(\underline{\rho}_3)} \, T(\underline{\rho}_2) T^*\left(\underline{\rho}_2 + \frac{\underline{\omega} L}{k}\right)$$

$$\cdot \exp\left\{ \psi(\underline{\rho}_2, \underline{\rho}_1) + \psi^*(\underline{\rho}_2 + \underline{\omega} L/k, \underline{\rho}_3) + i \frac{k}{x} \underline{\rho}_2 \cdot (\underline{\rho}_1 - \underline{\rho}_3) - i \frac{L}{x} \underline{\omega} \cdot \underline{\rho}_3 \right\} \tag{4}$$

The overbar on the quantity $u_o(\underline{\rho}_1) u_o^*(\underline{\rho}_3)$ indicates an average over source statistics, and arises because it has been tacitly assumed that the exposure time of the individual speckle pattern is much longer than the source coherence time, but of course, much shorter than the coherence time of the turbulent atmosphere.

For a spatially incoherent source one usually assumes $\overline{u_o(\underline{\rho}_1) u_o^*(\underline{\rho}_3)} = A_c I_s(\underline{\rho}_1) \delta(\underline{\rho}_1 - \underline{\rho}_3)$ where $\delta(\ldots)$ is the Dirac delta function and $I_s$ is the source irradiance distribution. However, for partially coherent sources this assumption is invalid, and instead we shall employ the quasi-stationary model used by Carter and Wolf,[10] and others. That is, we shall write

$$\overline{u_o(\underline{\rho}_1) u_o^*(\underline{\rho}_3)} \approx I_s\left(\frac{\underline{\rho}_1 + \underline{\rho}_3}{2}\right) g(\underline{\rho}_1 - \underline{\rho}_3) \tag{5}$$

where g is the correlation function of the source. In writing Eq. (5), it is tacitly assumed that the source coherence length is much smaller than the total size of the source.

We now use Eq. (5) in Eq. (4), define sum and difference coordinates $\underline{\xi} = \underline{\rho}_1 - \underline{\rho}_3$, $\underline{\eta} = (\underline{\rho}_1 + \underline{\rho}_3)/2$, etc., and then perform* the ensemble average over the turbulence statistics, assuming that the source lies wholly within an isoplanatic turbulence patch, so that $\psi(\underline{\rho}, \underline{\rho}') \simeq \psi(\rho, 0)$. The result is

$$\langle | I_f(\underline{\omega}) |^2 \rangle = H(\underline{\omega}) \left| \hat{I}_s \left( \underline{\omega} \frac{L}{x} \right) \right|^2 , \tag{6}$$

where

$$\hat{I}_s(\underline{\Omega}) = (2\pi)^{-2} \iint_{-\infty}^{\infty} d^2 v \, I_s(\underline{v}) \exp(-i \underline{\Omega} \cdot \underline{v}) . \tag{7}$$

Therefore, the modulus squared of the Fourier transform of the source distribution is linearly related to the measured quantity, $\langle | I_f |^2 \rangle$. Also in Eq. (6), we have defined

$$H(\underline{\omega}) = (\lambda x)^{-4} \iint_{-\infty}^{\infty} d^2\gamma \, K_o(\underline{\omega}, \underline{\gamma}) \exp \left\{ - D_1 \left( \underline{\omega} \frac{L}{k}, 0 \right) - D_1(\underline{\gamma}, 0) \right.$$
$$\left. + \frac{1}{2} D_1 \left( \underline{\gamma} - \underline{\omega} \frac{L}{k}, 0 \right) + \frac{1}{2} D_1 \left( \underline{\gamma} + \underline{\omega} \frac{L}{k}, 0 \right) \right\} , \tag{8}$$

where

$$K_o(\underline{\omega}, \underline{\gamma}) = \iint_{-\infty}^{\infty} d^2\tau \, T(\underline{\tau} + \underline{\gamma}/2) \, T^*(\underline{\tau} + \underline{\gamma}/2 + \underline{\omega} L/k)$$
$$\cdot T^*(\underline{\tau} - \underline{\gamma}/2) \, T(\underline{\tau} - \underline{\gamma}/2 + \underline{\omega} L/k) \, G(\underline{\tau}, \underline{\gamma}, \underline{\omega}) \, G^*(\underline{\tau}, -\underline{\gamma}, \underline{\omega}) , \tag{9}$$

---

*The derivation of the average,

$$\langle \exp [\psi(\underline{\rho}_2, \underline{\rho}_1) + \psi^*(\underline{\rho}_2 + \underline{\omega} L/k, \underline{\rho}_1) + \psi^*(\underline{\rho}_5, \underline{\rho}_4) + \psi(\underline{\rho}_5 + \underline{\omega} L/k, \underline{\rho}_4)] \rangle$$

over the statistics of the turbulence is given by Eq. (34) of R. Fante, "Some results on the imaging of incoherent sources through turbulence," J. Opt. Soc. Am. 66, 574-580 (1976).

and

$$G(\underline{\tau}, \underline{\gamma}, \underline{\omega}) = \iint_{-\infty}^{\infty} d^2\xi \; g(\underline{\xi}) \exp \left\{ i \frac{k}{x} \underline{\xi} \cdot (\underline{\tau} + \underline{\gamma}/2 + \underline{\omega} L/2k) \right\} . \tag{10}$$

The structure function $D_1$ of the atmospheric turbulence is given by

$$D_1(\underline{\beta}, 0) \simeq 2 \, |\beta/r_o|^{5/3} \quad , \tag{11}$$

and the atmospheric coherence length, $r_o$, is defined as

$$r_o = \left[ 1.46 \, k^2 \int_0^x \left( \frac{x - x'}{x} \right)^{5/3} C_n^2(x') \, dx' \right]^{-3/5} , \tag{12}$$

where $C_n^2(x')$ is the index-of-refraction structure constant of the atmosphere.

The result in Eq. (6) is of the same form as previous results obtained for a spatially incoherent source, but differs in the definition of $K_o(\underline{\omega}, \underline{\gamma})$. For a spatially incoherent source, so that $g(\underline{\rho}_1 - \underline{\rho}_3) = A_c \, \delta \, (\underline{\rho}_1 - \underline{\rho}_3)$ where $A_c$ is the coherence area of the source, we find $G = A_c$, and Eqs. (6)–(10) then reduce identically to the result of Korff[1] for the spatially incoherent source.

## 3. APPROXIMATE EVALUATION OF TRANSFER FUNCTION, $H(\underline{\omega})$

We would now like to investigate the effect of the partial coherence on the transfer function, $H(\underline{\omega})$. In order to do this, we will assume that the source coherence function $g(\underline{\xi})$ is given by

$$g(\underline{\xi}) = \exp \left( - \frac{\xi^2}{a^2} \right) , \tag{13}$$

where a is the source coherence length. Upon using this expression in Eq. (10), we find

$$G(\underline{\tau}, \underline{\gamma}, \underline{\omega}) = A_c \exp \left[ - \left( \frac{ka}{2x} \right)^2 |\underline{\tau} + \underline{\gamma}/2 + \underline{\omega} L/2k|^2 \right] , \tag{14}$$

where $A_c = \pi a^2$ is the coherence area of the source.

The next step is to substitute G into Eq. (9) and evaluate $K_o(\underline{\omega}, \underline{\gamma})$. For a realistic aperture, such as on a telescope, $T(\underline{\rho}_2) = 1$ for $|\underline{\rho}_2| \leq D$ and $T(\underline{\rho}_2) = 0$ for $|\rho_2| > D$. Unfortunately, the expression for K in Eq. (9) cannot be* evaluated in closed form when this expression is used for $T(\underline{\rho}_2)$. Therefore, we have been forced to approximate $T(\underline{\rho}_2)$ by the function $T(\underline{\rho}_2) = \exp(-\rho_2^2/D^2)$, with the realization that, because we are approximating the edge diffraction, our results will be valid only for radian spatial frequencies such that $|\underline{\omega}| < kD/L$. That is, setting $T(\underline{\rho}_2) = 0$ for $|\underline{\rho}_2| > D$ would give $H(\underline{\omega}) = 0$ for $|\underline{\omega}| > kD/L$; however, because we will approximate $T(\underline{\rho}_2)$ by $\exp(-\rho_2^2/D^2)$ we will obtain an expression for $H(\underline{\omega})$ which does not vanish for $|\underline{\omega}| > kD/L$. If the aforementioned approximation is used, along with Eq. (14), in Eq. (9) we find that

$$K_o(\underline{\omega}, \underline{\gamma}) = \frac{\frac{\pi}{4} A_c^2 \exp\left\{-\frac{\omega^2 L^2}{k^2 D^2} - \gamma^2\left(\frac{1}{\ell_c^2} + \frac{1}{D^2}\right)\right\}}{\left(\frac{1}{\ell_c^2} + \frac{1}{D^2}\right)} , \tag{15}$$

where $\ell_c = 2^{3/2} x/(ka) = (2/\pi A_c)^{1/2} \lambda x$ is the coherence length, as measured at the receiving aperture, of the field radiated by a coherence patch of radius a on the source.

We next investigate the behavior of the function

$$C(\underline{\omega}, \underline{\gamma}) = \exp\{-D_1(\underline{\omega} L/k, 0) - D_1(\underline{\gamma}, 0) + \tfrac{1}{2} D_1(\underline{\gamma} - \underline{\omega} L/k, 0) + \tfrac{1}{2} D_1(\underline{\gamma} + \underline{\omega} L/k, 0)\} . \tag{16}$$

It is found that for $|\underline{\omega}| \ll kr_o/L$, this function can be approximated by unity. However, the more interesting case occurs** when $|\underline{\omega}| \gg kr_o/L$. In this limit, $C(\underline{\omega}, \underline{\gamma})$ can be approximated by

$$C(\underline{\omega}, \underline{\gamma}) \simeq \exp\{-2|\underline{\gamma}/r_o|^{5/3}\} . \tag{17}$$

The result in Eq. (17) is inconvenient for analytical purposes; consequently, we shall approximate the exponent in Eq. (17) by a quadratic function, so that

*The evaluation can be performed in the limit when $G(\xi) = A_c$.

**We are assuming that $r_o < D$. For $r_o \gg D$, the atmospheric turbulence has an insignificant effect on the imaging.

$$C(\underline{\omega}, \underline{\gamma}) \simeq \exp \left\{ - 2.296 \, |\underline{\gamma}/r_o|^2 \right\} \quad . \tag{18}$$

The function in Eq. (18) is such that C equals $e^{-1}$ at the same value of $|\underline{\gamma}|$ as does the function in Eq. (17).

If we now use Eqs. (15) and (18) in Eq. (8), we can calculate the transfer function $H(\underline{\omega})$. For values of $|\underline{\omega}|$ such that $|\underline{\omega}| \gg kr_o/L$, the result is*

$$H(\underline{\omega}) = \frac{K \exp(-\omega^2 L^2 / k^2 D^2)}{\left[1 + \left(\frac{D}{\ell_c}\right)^2\right] \left[1 + \left(\frac{D}{\ell_c}\right)^2 + 2.296 \left(\frac{D}{r_o}\right)^2\right]} \quad , \tag{19}$$

where

$$K = 2.467 \left(\frac{D}{x}\right)^4 \left(\frac{A_c^2}{\lambda^4}\right) = \left(\frac{D}{\ell_c}\right)^4 \quad . \tag{20}$$

Because of our approximation for $T(\underline{\rho}_2)$, the result in Eq. (19) can be applied to realistic apertures only for spatial frequencies such that $|\underline{\omega}| < kD/L$. In the limit when $r_o \ll D$ and the source is incoherent, so that** $\ell_c \to \infty$, Eq. (19) reduces to the high-spatial frequency approximation obtained by Korff.[1]

For values of $|\underline{\omega}|$ such that $|\underline{\omega}| \ll kr_o/L$ the evaluation of Eq. (8) gives†

$$H(\underline{\omega}) = \frac{K}{\left[1 + \left(\frac{D}{\ell_c}\right)^2\right]^2} \quad . \tag{21}$$

A qualitative plot of $H(\omega)$ is shown in Figure 2, for the limits when $r_o \ll D$ and $r_o \gg D$. In this latter case, of course, the turbulent atmosphere does not affect the imaging.

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*Note that $H(\omega)$ is not normalized to unity at $\omega = 0$, as is often done.

**Strictly speaking $\ell_c$ never really becomes infinite, because $a \gtrsim \lambda$. Therefore, $\ell_c \to 2^{1/2} \pi^{-1} x$.

†Because $r_o$ is assumed to be less than D, the assumption that $|\omega| \ll kr_o/L$ implies that $|\omega L/kD| \ll 1$. Consequently, in Eq. (21) we have approximated $\exp(-\omega^2 L^2/k^2 D^2)$ by unity.

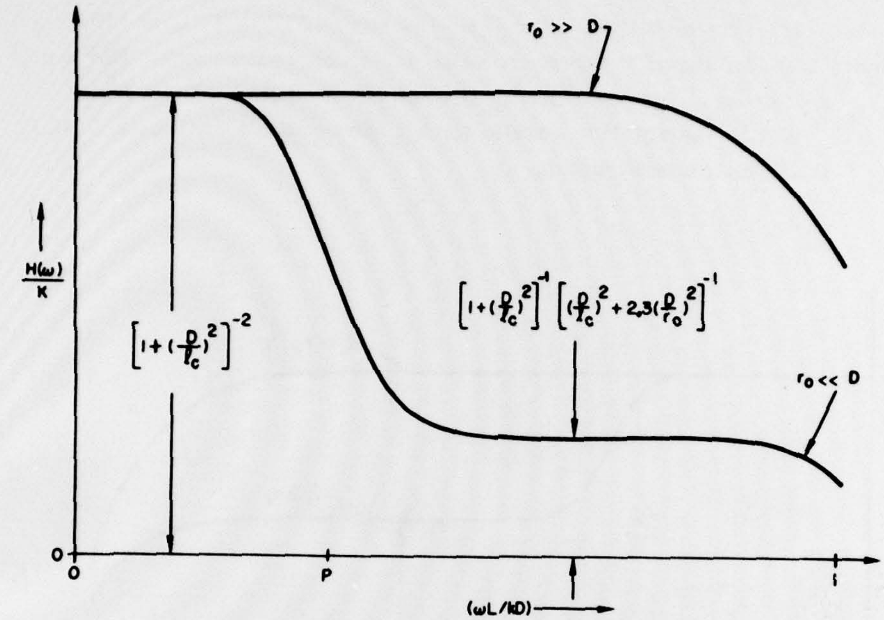


Figure 2. Qualitative Plot of $H(\underline{\omega})$. The point P corresponds to $|\underline{\omega}| \simeq kr_o/L$

## 4. DISCUSSION

From Figure 2, we observe that the amplitude of $H(\underline{\omega})$ at small spatial frequencies is determined by the size of $D/\ell_c$. For $\ell_c \gg D$ we get the usual Labeyrie result, but when $\ell_c \ll D$ the result in Eq. (21) becomes $H(\underline{\omega}) \rightarrow 1$.

It is also interesting to observe the nature of the high-spatial-frequency plateau in $H(\underline{\omega})$ for the case when $r_o \ll D$. We will not discuss the case when $r_o \gg D$ because in that limit atmospheric turbulence does not affect the image. If $r_o \ll D$ but $\ell_c \gg D$, we get the previous result of Korff; however, if $r_o \ll D$ and $\ell_c \ll D$ we find, for $|\underline{\omega}| > kr_o/L$

$$H(\omega) \simeq \frac{\exp\left(-\omega^2 L^2/k^2 D^2\right)}{1 + 2.296\left(\frac{\ell_c}{r_o}\right)^2} \quad . \tag{22}$$

From Eq. (22) we see that if $\ell_c \ll r_o$, $H(\underline{\omega}) \to \exp(-\omega^2 L^2/k^2 D^2)$, and the high-frequency transfer function approaches that for a coherent source. However, when $\ell_c$ is of order or greater than $r_o$ the amplitude of $H(\underline{\omega})$ differs from $\exp(-\omega^2 L^2/k^2 D^2)$. Qualitative results for $H(\underline{\omega})$ in the limit when both $\ell_c \ll D$ and $r_o \ll D$ are shown in Figure 3.

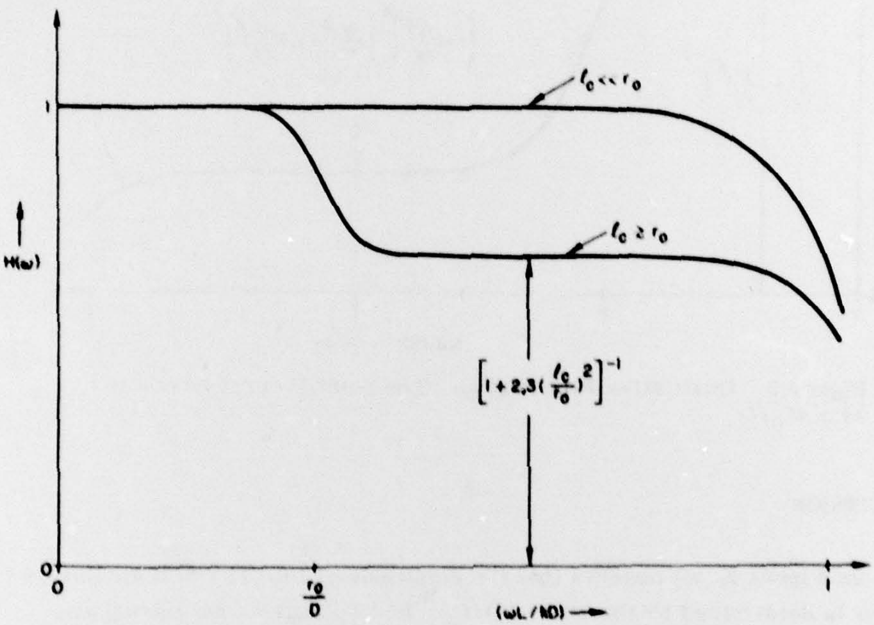


Figure 3. Qualitative Plot of $H(\underline{\omega})$ for the Case When Both $\ell_c \ll D$ and $r_o \ll D$

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